Almost the supersymmetric Standard Model from intersecting D6-branes on the $\mathbb{Z}_6'$ orientifold

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Abstract

Intersecting stacks of $\mathcal{N}=1$ supersymmetric fractional branes on the $\mathbb{Z}_6'$ orientifold may be used to construct the supersymmetric Standard Model. If a, b are the stacks that generate the $SU(3)_{\mathrm{colour}}$ and $SU(2)_L$ gauge particles, then, in order to obtain *just* the chiral spectrum of the (supersymmetric) Standard Model (with non-zero Yukawa couplings to the Higgs mutiplets), it is necessary that the number of intersections $a \cap b$ of the stacks a and b, and the number of intersections $a \cap b'$ of a with the orientifold image b' of b satisfy $(a \cap b, a \cap b') = (2,1)$ or $(1,2)$. It is also necessary that there is no matter in symmetric representations of the gauge group, and not too much matter in antisymmetric representations, on either stack. Fractional branes having all of these properties may be constructed on the $\mathbb{Z}_6'$ orientifold. We construct a (four-stack) model with two further stacks, each with just a single brane, which has precisely the matter spectrum of the supersymmetric Standard Model, including a single pair of Higgs doublets. However, the gauge group is $SU(3)_{\mathrm{colour}} \times SU(2)_L \times U(1)_Y \times U(1)_H$. Only the Higgs doublets are charged with respect to $U(1)_H$.

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An attractive, bottom-up approach to constructing the Standard Model is to use intersecting D6-branes [1]. In these models one starts with two stacks, a and b with $N_a = 3$ and $N_b = 2$, of D6-branes wrapping the three large spatial dimensions plus 3-cycles of the six-dimensional internal space (typically a torus T^6 or a Calabi-Yau 3-fold) on which the theory is compactified. These generate the gauge group $U(3) \times U(2) \supset SU(3)_c \times SU(2)_L$, and the non-abelian component of the standard model gauge group is immediately assured. Further, (four-dimensional) fermions in bifundamental representations $(\mathbf{N}_a, \overline{\mathbf{N}}_b) = (\mathbf{3}, \overline{\mathbf{2}})$ of the gauge group can arise at the multiple intersections of the two stacks. These are precisely the representations needed for the quark doublets Q_L of the Standard Model, and indeed an attractive model having just the spectrum of the Standard Model has been constructed [2]. The D6-branes wrap 3-cycles of an orientifold T^6/Ω, where Ω is the world-sheet parity operator. The advantage and, indeed, the necessity of using an orientifold stems from the fact that for every stack $a, b, ...$ there is an orientifold image $a', b', ...$ At intersections of a and b there are chiral fermions in the $(\mathbf{3}, \overline{\mathbf{2}})$ representation of $U(3) \times U(2)$, where the $\mathbf{3}$ has charge $Q_a = +1$ with respect to the $U(1)_a$ in $U(3) = SU(3)_{\text{colour}} \times U(1)_a$, and the $\overline{\mathbf{2}}$ has charge $Q_b = -1$ with respect to the $U(1)_b$ in $U(2) = SU(2)_L \times U(1)_b$. However, at intersections of a and b' there are chiral fermions in the $(\mathbf{3}, \mathbf{2})$ representation, where the $\mathbf{2}$ has $U(1)_b$ charge $Q_b = +1$. In the model of [2], the number of intersections $a \cap b$ of the stack a with b is 2, and the number of intersections $a \cap b'$ of the stack a with b' is 1. Thus, as required for the Standard Model, there are 3 quark doublets. (They have the same weak hypercharge Y provided that Q_b does not contribute to Y.) These have net $U(1)_a$ charge $Q_a = 6$, and net $U(1)_b$ charge $Q_b = -3$. Tadpole cancellation requires that overall both charges, sum to zero, so further fermions are essential, and indeed required by the Standard Model. 6 quark-singlet states u^c_L and d^c_L belonging to the $(\mathbf{1}, \overline{\mathbf{3}})$ representation of $U(1) \times U(3)$, having a total of $Q_a = -6$ are sufficient to ensure overall cancellation of Q_a, and these arise from the intersections of a with other stacks $c, d, ...$ having just a single D6-brane. Similarly, 3 lepton doublets L, belonging to the $(\mathbf{2}, \overline{\mathbf{1}})$ representation of $U(2) \times U(1)$, having a total $U(1)_b$ charge of $Q_b = 3$, are sufficient to ensure overall cancellation of Q_b, and these arise from the intersections of b with other stacks having just a single D6-brane. In contrast, had we not used an orientifold, the requirement of 3 quark doublets would necessitate having the number of intersections $a \cap b = 3$. This makes no difference to the charge $Q_a = 6$ carried by the quark doublets, but instead the $U(1)_b$ charge carried by the quark doublets is $Q_b = -9$, which cannot be cancelled by just 3 lepton doublets L. Consequently, additional vector-like fermions are unavoidable unless the orientifold projection is available. This is why the orientifold is essential if we are to get just the matter content of the Standard Model or of the MSSM.

Actually, an orientifold can allow essentially the standard-model spectrum without vector-like matter even when $a \cap b = 3$ and $a \cap b' = 0$ [3]. This is because in orientifold models it is also possible to get chiral matter in the symmetric and/or antisymmetric representation of the relevant gauge group from open strings stretched between a stack and its orientifold image. Both representations have charge $Q = 2$ with respect to the relevant $U(1)$. The antisymmetric (singlet) representation of $U(2)$ can describe a neutrino singlet state ν^c_L (since Q_b does not contribute to Y), and 3 copies contribute $Q_b = 6$ units of $U(1)_b$ charge. If there are also 3 lepton doublets L belonging to the bifundamental representation $(\mathbf{2}, \overline{\mathbf{1}})$ representation of $U(2) \times U(1)$, each contributing $Q_b = 1$ as above, then the total contribution is $Q_b = 9$ which **can** be cancelled by 3 quark doublets Q_L in the $(\mathbf{3}, \overline{\mathbf{2}})$ representation of $U(3) \times U(2)$. Thus, orientifold models can allow the standard-model spectrum plus 3 neutrino singlet states even when $(a \cap b, a \cap b') = (3, 0)$.

Non-supersymmetric intersecting-brane models lead to flavour-changing neutral-current (FCNC) processes that can only be suppressed to levels consistent with the current bounds by making the string scale rather high, of order 10^4 TeV, which in turn leads to fine-tuning problems [4]. Further, in non-supersymmetric theories, such as these, the cancellation of Ramond-Ramond (RR) tadpoles does not ensure Neveu Schwarz-Neveu Schwarz (NSNS) tadpole cancellation. Thus a particular consequence of the non-cancellation is that the complex structure moduli are unstable [5]. One way to stabilise these moduli is for the D-branes to wrap an orbifold T^6/P, where P is a "point group" acting on T^6, rather than a torus T^6. The FCNC problem can be solved and the complex structure moduli stabilised when the theory is supersymmetric. First, a supersymmetric theory is not obliged to have the low string scale that led to problematic FCNCs induced by string instantons. Second, in a supersymmetric theory, RR tadpole cancellation ensures cancellation of the NSNS tadpoles [6, 7]. An orientifold is then constructed by quotienting the orbifold with the world-sheet parity operator Ω.

In this paper we shall be concerned with the orientifold having point group $P = \mathbb{Z}_6'$. We showed in a previous paper [8] that this *does* have (fractional) supersymmetric D6-branes a and b with intersection numbers $(a \cap b, a \cap b') = (1,2)$ or $(2,1)$, which might be used to construct the supersymmetric Standard Model having just the requisite standard-model matter content.

The torus T^6 factorises into three 2-tori $T_1^2 \times T_2^2 \times T_3^2$, with T_k^2 $(k = 1,2,3)$ parametrised by the complex coordinate z_k. The action of the generator θ of the point group $\mathbb{Z}_6'$ on the coordinates z_k is given by

$$\theta z_k = e^{2\pi i v_k} z_k \tag{1}$$

where

$$(v_1, v_2, v_3) = \frac{1}{6}(1, 2, -3) \tag{2}$$

Since the point group action must be an automorphism of the lattice, we take $T_{1,2}^2$ to be $SU(3)$ lattices, so that their complex structure moduli satisfy $U_1 = e^{i\pi/3} = U_2$, whereas the lattice for and T_3^2 and hence its complex structure U_3 is arbitrary. The anti-holomprphic embedding $\mathcal{R}$ of the world-sheet parity operator Ω acts as complex conjugation on the coordinates z_k

$$\mathcal{R} z_k = \overline{z}_k \quad (k = 1,2,3) \tag{3}$$

Requiring that this too is an automorphism of the lattice constrains the orientation of each torus T_k^2 relative to the Re z_k axis. Each torus must be in one of two configurations, denoted **A** and **B**, defined in reference [8]. This fixes the real part of the complex structure Re $U_3 = 0, 1/2$ respectively, but the imaginary part Im U_3 remains *a priori* arbitrary. In this paper, we shall only be concerned with the **ABA** lattice.

The fractional branes κ with which we are concerned have the general form

$$\kappa = \frac{1}{2}\left(\Pi_\kappa^{\text{bulk}} + \Pi_\kappa^{\text{ex}}\right) \tag{4}$$

where

$$\Pi_\kappa^{\text{bulk}} = \sum_{p=1,3,4,6} A_p^\kappa \rho_p \tag{5}$$

is an (untwisted) invariant 3-cycle, and

$$\Pi_\kappa^{\text{ex}} = \sum_{j=1,4,5,6} (\alpha_j^\kappa \epsilon_j + \tilde{\alpha}_j^\kappa \tilde{\epsilon}_j) \tag{6}$$

is an exceptional 3-cycle associated with the θ^3-twisted sector. It consists of a collapsed 2-cycle at a θ^3 fixed point in $T_1^2 \times T_3^2$ times a 1-cycle in the (θ^3-invariant plane) T_2^2. The 4 basis invariant 3-cycles ρ_p, $(p = 1,3,4,6)$ and the 8 basis exceptional cycles ϵ_j and $\tilde{\epsilon}_j$, $(j = 1,4,5,6)$ are defined in reference [8]. Their non-zero intersection numbers are

$$\rho_1 \cap \rho_4 = -4, \qquad \rho_1 \cap \rho_6 = 2 \tag{7}$$

$$\rho_3 \cap \rho_4 = 2, \qquad \rho_3 \cap \rho_6 = -4 \tag{8}$$

and

$$\epsilon_j \cap \tilde{\epsilon}_k = -2\delta_{jk} \tag{9}$$

The wrapping numbers (n_k^a, m_k^a) of the basis 1-cycles (π_{2k-1}, π_{2k}) of T_k^2 for the $U(3)$ stack a are given by

$$(n_1^a, m_1^a; n_2^a, m_2^a, n_3^a, m,_3^a) = (1, -2; -1, 0; 1, -2) \tag{10}$$

Using the formulae given in [8],

$$A_1 = (n_1 n_2 + n_1 m_2 + m_1 n_2) n_3 \tag{11}$$

$$A_3 = (m_1 m_2 + n_1 m_2 + m_1 n_2) n_3 \tag{12}$$

$$A_4 = (n_1 n_2 + n_1 m_2 + m_1 n_2) m_3 \tag{13}$$

$$A_6 = (m_1 m_2 + n_1 m_2 + m_1 n_2) m_3 \tag{14}$$

we can compute the bulk coefficients A^a_p for $\Pi^{\rm bulk}_a$. Then the fractional brane a has

$$\begin{aligned}\Pi^{\rm bulk}_a &= \rho_1 + 2\rho_3 - 2\rho_4 - 4\rho_6 &(15)\\ \Pi^{\rm ex}_a &= (-1)^{\tau^a_0}\left(2[\epsilon_1 + (-1)^{\tau^a_2}\epsilon_4] + [\tilde{\epsilon}_1 + (-1)^{\tau^a_2}\tilde{\epsilon}_4]\right) &(16)\end{aligned}$$

On the **ABA** lattice, the orientifold O6-plane is

$$\Pi_{O6} = 2\rho_1 + \rho_3 - 3\rho_6 \tag{17}$$

and the orientifold images of $\Pi^{\rm bulk}_a$ and $\Pi^{\rm ex}_a$ are

$$\begin{aligned}\Pi^{\rm bulk\prime}_a &= \rho_1 - \rho_3 + 2\rho_4 - 2\rho_6 &(18)\\ \Pi^{\rm ex\prime}_a &= = (-1)^{\tau^a_0}\left(-[\epsilon_1 + (-1)^{\tau^a_2}\epsilon_4] + [\tilde{\epsilon}_1 + (-1)^{\tau^a_2}\tilde{\epsilon}_4]\right) &(19)\end{aligned}$$

Then

$$\begin{aligned}a \cap \Pi_{O6} &= \frac{1}{2}\Pi^{\rm bulk}_a \cap \Pi_{O6} = 3 &(20)\\ &= a \cap a' &(21)\end{aligned}$$

from which it follows, as required, that there are no symmetric representations $\mathbf{S}_a = \mathbf{6}$ on the stack a. On this lattice, supersymmetry requires that the bulk coefficients A^a_p satisfy

$$\begin{aligned}&\sqrt{3}A^a_1 + (A^a_4 - 2A^a_6)\,{\rm Im}\,U_3 > 0 &(22)\\ &-A^a_1 + 2A^a_3 + A^a_4\sqrt{3}\,{\rm Im}\,U_3 = 0 &(23)\end{aligned}$$

where U_3 is the complex structure on T^2_3. Using the bulk coefficients for a given in (15), it follows that

$${\rm Im}\,U_3 = \frac{\sqrt{3}}{2} \tag{24}$$

Likewise, the wrapping numbers (n^b_k, m^b_k) for the $U(2)$ stack b are given by

$$(n^b_1, m^b_1; n^b_2, m^b_2, n^b_3, m^b_{,3}) = (0, 1; 0, -1; 0, 1) \tag{25}$$

and the fractional brane b has

$$\begin{aligned}\Pi^{\rm bulk}_b &= -\rho_6 = \Pi^{\rm bulk\prime}_b &(26)\\ \Pi^{\rm ex}_b &= = (-1)^{\tau^b_0+1}[\epsilon_1 + (-1)^{\tau^b_2}\epsilon_5] = -\Pi^{\rm ex\prime}_b &(27)\end{aligned}$$

It follows that b too is supersymmetric and that

$$b \cap \Pi_{O6} = 0 = b \cap b' \tag{28}$$

so, again, there are no symmetric representations $\mathbf{S}_b = \mathbf{3}$ on the stack b. Then, as required,

$$(a \cap b, a \cap b') = (2, 1) \quad \text{or} \quad (1, 2) \tag{29}$$

the former occurring when $\tau^a_0 = \tau^b_0 \bmod 2$ and the latter when $\tau^a_0 \neq \tau^b_0 \bmod 2$.

The weak hypercharge Y is a linear combination

$$Y = \sum_\kappa y_\kappa Q_\kappa \tag{30}$$

of the charges Q_κ associated with the $U(1)_\kappa$ group for each stack κ. As explained above, the intersections $a \cap b$ give chiral supermultiplets in the $(\mathbf{3}, \overline{\mathbf{2}})$ representation of $U(3) \times U(2)$ and $a \cap b'$ give chiral supermultiplets in the $(\mathbf{3},\mathbf{2})$ representation. Since both of these give quark doublets Q_L having $Y = 1/6$, we infer that

$$\begin{aligned} y_a &= \frac{1}{6} && (31)\\ y_b &= 0 && (32)\end{aligned}$$

(As noted earlier, the stack Q_b makes no contribution to Y.) Using (20) we find that the number of antisymmetric representations on the stack a is

$$\#(\mathbf{A}_a) = \frac{1}{2}(a \cap a' + a \cap \Pi_{O6}) = 3 \tag{33}$$

Since the antisymmetric part $\mathbf{A}_a = (\mathbf{3} \times \mathbf{3})_{\text{antisymm}} = \overline{\mathbf{3}}$ for the group $SU(3)$, and since $\mathbf{3}$ has $Q_a = 1$, it follows that there are 3 $(\overline{\mathbf{3}},\mathbf{1})$ representations of $SU(3)_{\text{colour}} \times SU(2)_L$ having $Y = 1/3$. Thus, there are 3 d^c_L quark-singlet states on the stack a. Similarly, from (28), it follows that

$$\#(\mathbf{A}_b) = 0 \tag{34}$$

so that there are no lepton-singlet ν^c_L states on the stack b. To complete the standard-model spectrum, it is therefore necessary to add further stacks, $c, d, e, \ldots$ each having just a single (fractional) brane $N_{c,d,e\ldots} = 1$. In principle, it might be possible to obtain the full spectrum with the addition of just one further stack c with $y_c = 1/2$. The remaining 3 quark-singlet states u^c_L having weak hypercharge $Y = -2/3$ might arise at the intersections of a and c, the 3 lepton and 2 Higgs doublets with $Y = \pm 1/2$ from intersections of b with c, with the 3 charged lepton-singlet states ℓ^c_L having $Y = 1$ possibly arising as symmetric representations on c. However, we have not so far been able to find such an example. The example that we shall present has two further $U(1)$ stacks.

The wrapping numbers (n^c_k, m^c_k) for the first of these stacks are

$$(n^c_1, m^c_1; n^c_2, m^c_2, n^c_3, m^c_{,3}) = (1,0;1,0;3,2) \tag{35}$$

and the fractional brane c has

$$\begin{aligned} \Pi^{\text{bulk}}_c &= 3\rho_1 + 2\rho_4 && (36)\\ \Pi^{\text{ex}}_c &= (-1)^{\tau^c_0+1}\left(2[\epsilon_1 + (-1)^{\tau^c_2}\epsilon_4] + [\tilde{\epsilon}_1 + (-1)^{\tau^c_2}\tilde{\epsilon}_4]\right) && (37)\end{aligned}$$

Then c is supersymmetric and

$$\begin{aligned} \Pi^{\text{bulk}\prime}_c &= 3\rho_1 + 3\rho_3 - 2\rho_4 - 2\rho_6 && (38)\\ \Pi^{\text{ex}\prime}_c &= = (-1)^{\tau^c_0+1}\left(-[\epsilon_1 + (-1)^{\tau^c_2}\epsilon_4] + [\tilde{\epsilon}_1 + (-1)^{\tau^c_2}\tilde{\epsilon}_4]\right) && (39)\end{aligned}$$

Taking

$$\begin{aligned} \tau^a_0 + \tau^c_0 &= 0 \bmod 2 && (40)\\ \tau^a_2 + \tau^c_2 &= 0 \bmod 2 && (41)\end{aligned}$$

and using (15) and (16) then gives

$$(a \cap c, a \cap c') = (0, -3) \tag{42}$$

Hence, if

$$y_c = \frac{1}{2} \tag{43}$$

we get just the required 3 quark-singlet states u^c_L with weak hypercharge $Y = -2/3$. Further, the number $\#(\mathbf{S}_c)$ of symmetric representations on the stack c is

$$\begin{aligned} \#(\mathbf{S}_c) &= \frac{1}{2}(c \cap c' - c \cap \Pi_{O6}) && (44)\\ &= 3 && (45)\end{aligned}$$

so that we get just the required 3 charged lepton-singlet states ℓ^c_L having weak hypercharge $Y = 1$. Similarly, using (26) and (27), we find that

$$(b \cap c, b \cap c') = (2, -1) \quad \text{or} \quad (1, -2) \tag{46}$$

the former occurring when $\tau^b_0 + \tau^c_0 = 0 \bmod 2$ and the latter when it is $1 \bmod 2$. Either way, this gives 3 lepton (or Higgs) doublets each having weak hypercharge $Y = -1/2$. Evidently a further stack is required to generate a pair of doublets with opposite weak hypercharges $Y = 1/2$ and $Y = -1/2$.

The wrapping numbers (n^d_k, m^d_k) for the second of the $U(1)$ stacks are

$$(n^d_1, m^d_1; n^d_2, m^d_2, n^d_3, m,^d_3) = (1, 0; 1, 1; 1, 0) \tag{47}$$

and the fractional brane d has

$$\begin{aligned} \Pi^{\text{bulk}}_d &= 2\rho_1 + \rho_3 = \Pi^{\text{bulk}'}_d & (48) \\ \Pi^{\text{ex}}_d &= (-1)^{\tau^d_0+1} \left([\epsilon_1 + (-1)^{\tau^d_2} \epsilon_4] + 2[\tilde{\epsilon}_1 + (-1)^{\tau^d_2} \tilde{\epsilon}_4] \right) = \Pi^{\text{ex}'}_d & (49) \end{aligned}$$

Then d is supersymmetric and

$$d \cap d' = 0 = d \cap \Pi_{O6} \tag{50}$$

so that there are no (lepton-singlet) states arising as the symmetric representation $\mathbf{S}_d$ on the stack d. Also, taking

$$\begin{aligned} \tau^a_0 + \tau^d_0 &= 0 \bmod 2 & (51) \\ \tau^a_2 + \tau^d_2 &= 0 \bmod 2 & (52) \end{aligned}$$

and using (15) and (16) then gives

$$(a \cap d, a \cap d') = (0, 0) \tag{53}$$

Thus there are no (unwanted) quark-singlet states arising at the intersections of a with d and d'. Using (26) and (27), we find that

$$(b \cap d, b \cap d') = (1, 1) \quad \text{or} \quad (-1, -1) \tag{54}$$

the former occurring when $\tau^b_0 + \tau^d_0 = 0 \bmod 2$ and the latter when it is $1 \bmod 2$. Either way, a pair of (Higgs) doublets, having the required opposite weak hypercharges occurs, provided that

$$y_d = \pm\frac{1}{2} \tag{55}$$

Finally, using (40),(41),(51) and (52), we find that

$$(c \cap d, c \cap d') = (0, 0) \tag{56}$$

so that there are no charged or neutral lepton singlets at the intersections of c with d or d'.

In total, we have *just* the spectrum of the supersymmetric Standard Model, with a single pair of Higgs doublets, and no neutrino singlet states ν^c_L. For this to be a consistent string theory realisation of the Standard Model it is necessary that there is overall cancellation of the RR tadpoles, and this in turn requires that the overall homology class of the D6-branes and O6-plane must vanish:

$$\sum_\kappa N_\kappa(\kappa + \kappa') = 4\Pi_{O6} \tag{57}$$

The sum is over all four stacks $\kappa = a, b, c, d$ and Π_{O6} is given in (17). Note that the left-hand side has contributions from both bulk and exceptional D6-branes, whereas the right-hand side has only the former. Both bulk and exceptional parts are required to cancel separately, and it is easy to verify that this is the case.

In the first instance the gauge group derived from these stacks is

$$\begin{aligned} G &= U(3)_a \times U(2)_b \times U(1)_c \times U(1)_d & (58) \\ &= SU(3)_{\text{colour}} \times SU(2)_L \times U(1)_Y \times U(1)^3 & (59) \end{aligned}$$

The last three $U(1)$s are of course unwanted. The gauge boson of any anomalous $U(1)$ gauge group acquires a mass of order the string scale, $O(10^{17})$ GeV in a supersymmetric theory, and the $U(1)$ survives only as a global symmetry. This is what happens in the case of $U(1)_a$, which is the $U(1)$ associated with (3 times the) baryon number. However, there remains the possibility that the gauge boson of a non-anomalous $U(1)$ symmetry remains massless (on the string scale) and might therefore be observable in low-energy experiments. The $U(1)$ gauge boson associated with a general linear combination of the $U(1)$ charges Q_κ

$$X = \sum_\kappa x_\kappa Q_\kappa \tag{60}$$

whether anomalous or non-anomalous, does *not* acquire a mass via the Green-Schwarz mechanism provided that

$$\sum_\kappa x_\kappa N_\kappa (\kappa - \kappa') = 0 \tag{61}$$

In the case under consideration, it is easy to verify that $U(1)_Y$ remains massless, as required, but so too does $U(1)_d$. Thus we have an unwanted further $U(1)$ factor in the gauge group. However, the only matter which has $Q_d \neq 0$ is the pair of doublets in (54) and it is attractive to identify these with the pair of Higgs doublets, H and $\overline{H}$. Since, Higgs particles have not yet been observed, this scenario has not yet been falsified.

In summary, we have found an attractive and economical realisation of the supersymmetric Standard Model using just four stacks of D6-branes on the **ABA** lattice. It has the correct matter content, with no neutrino-singlet states. There is one additional $U(1)$ factor in the gauge group, which may only interact with Higgs doublets. Although the complex structure moduli are stabilised in this model, there of course remain unstabilised Kähler and dilaton moduli. In principle, these may be stabilised using background fluxes, perhaps via the "rigid corset" proposed in [9]. In any case, fluxes are presumably needed to break the $\mathcal{N} = 1$ supersymmetry of the spectrum, as well as the $\mathcal{N} = 2$ supersymmetry of the gauge supermultiplets. These matters will be discussed elsewhere. Likewise, the results of a systematic investigation of other possible lattices for the $\mathbb{Z}_6'$ orientifold will be presented in a separate paper [10].

References

[1] For a review, see D. Lüst, Intersecting brane worlds: A path to the standard model?, Class. Quant. Grav. **21** (2004) S1399 [hep-th/0401156].

[2] L. E. Ibáñez, F. Marchesano and R. Rabadán, Getting just the standard model at intersecting branes, JHEP **0111** (2001) 002 [hep-th/0105155].

[3] R. Blumenhagen, B. Kors, D. Lust and T. Ott, "The standard model from stable intersecting brane world orbifolds," Nucl. Phys. B **616** (2001) 3 [arXiv:hep-th/0107138].

[4] S. A. Abel, O. Lebedev and J. Santiago, Flavour in intersecting brane models and bounds on the string scale, Nucl. Phys. B **696** (2004) 141 [hep-ph/0312157].

[5] R. Blumenhagen, B. Körs, D. Lüst and T. Ott, Intersecting brane worlds on tori and orbifolds, Fortsch. Phys. **50** (2002) 843 [hep-th/0112015].

[6] M. Cvetič, G. Shiu and A. M. Uranga, Three-family supersymmetric standard like models from intersecting brane worlds, Phys. Rev. Lett. **87** (2001) 201801 [hep-th/0107143].

[7] M. Cvetič, G. Shiu and A. M. Uranga, Chiral four-dimensional N = 1 supersymmetric type IIA orientifolds from intersecting D6-branes, Nucl. Phys. B **615** (2001) 3 [hep-th/0107166].

[8] D. Bailin and A. Love, "Towards the supersymmetric standard model from intersecting D6-branes on the Z'(6) orientifold," Nucl. Phys. B **755** (2006) 79 [arXiv:hep-th/0603172].

[9] P. G. Camara, A. Font and L. E. Ibanez, JHEP **0509** (2005) 013 [arXiv:hep-th/0506066].

[10] D. Bailin and A. Love, in preparation.